

# Stable Models of Super-acceleration

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We discuss an instability in a large class of models where dark energy is coupled to matter. In these models the mass of the scalar field is much larger than the expansion rate of the universe. We find models in which this instability is absent, and show that these models generically predict an *apparent* equation of state for dark energy smaller than -1, i.e., super-acceleration. These models have no acausal behavior or ghosts.

## MOTIVATION

Observations of distant Type Ia supernovae [1, 2] and the cosmic microwave background [3] together strongly prefer an accelerated expansion of the universe in the recent past. In the standard cosmological model this is accommodated by introducing “dark energy”, a component which has a significantly negative pressure causing the expansion of the universe to accelerate.

In the standard cosmological model, dark energy is completely decoupled from the rest of the matter in the universe except for its gravitational effects. It is interesting to consider more general models in which the dark matter and dark energy have a coupling. Such models could have new nontrivial signatures in cosmology and structure formation.

One simple class of such models is a model in which the vacuum energy density depends on the matter density. We shall consider a class of these models in which the dark energy responds to changes in the matter density on a time scale shorter than the expansion time scale. For example, one can consider models with scalar field dark energy coupled to matter (e.g., [4, 5, 6, 7, 8, 9, 10]), in which the mass of the scalar field is much larger than the expansion rate (for example, the MaVaN scenario [11]).

As we show below, these models generically suffer from an instability which we label AZK-instability. The AZK-instability was pointed out in the context of mass-varying neutrinos (MaVaN) [12]. A similar effect was identified in the context of unified dark energy models [13]. This instability can also occur in models of dark energy coupled to matter, such as the MaVaN scenario [11], the Chameleon dark energy scenario [14] and the Cardassian expansion scenario [15]. Not all models in the above scenarios are necessarily unstable (for example, [16, 17, 18]). This will become clear when we discuss the instability.

In this paper, we will construct a large class of models in which this instability is avoided. We find that these models generically predict an apparent equation of state (pressure over energy density)  $w_{\text{DE}}$  which is less than -1 (such a phase is labeled super-acceleration [19]). That is, a model of interacting dark energy can be incorrectly interpreted as a theory with super-acceleration if the in-

teractions are not taken into account.

For example, the coupling of dark energy to matter could be such that the total matter density decreases more slowly than  $1/a^3$  (where  $a$  is the scale factor of the universe). When we interpret observations in such a universe with a canonical matter density term (that decreases with expansion as  $1/a^3$ ) and dark energy, we would infer an equation of state for dark energy more negative than it truly is [20, 21]. There is no physical reason why this inferred equation of state cannot be below -1.

This is particularly interesting because current data seem to favor a dark energy density which is almost constant or even increasing with time [22, 23, 24, 25, 26, 27, 28, 29, 30, 31]. and exciting results can be expected in the future [32, 33, 34, 35]. SNIa observations currently favor a phase of super-acceleration. Future SNIa and CMB observations have the potential to detect super-acceleration [19]. No other combination has been shown to robustly detect the signature of super-acceleration, although combining SNIa and baryon oscillation [30] or weak lensing data set seem promising. Note that a measurement of just the average equation of state [36] is not sufficient for this purpose [37]. This was made explicit recently [38] using a simple single scalar field model.

Scalar field models with canonical kinetic terms always produce  $w_{\text{DE}} > -1$ . Effective models with the opposite sign kinetic term [22, 39] imply  $w_{\text{DE}} < -1$  but are unstable [40] unless more than one scalar field [41, 42, 43, 44, 45] or quantum effects [46] are considered. Models with higher derivative terms or scalar-tensor theories can give rise to an apparent  $w_{\text{DE}} < -1$  [47], but are constrained [48, 49, 50]. Interpreting an alternative gravity theory in the context of 4-d GR can also lead to super-acceleration [51, 52, 53, 54, 55, 56]. Some Cardassian models may have  $w_{\text{DE}} < -1$  [57, 58, 59] while still satisfying the dominant energy condition. Another possible way to get super-acceleration with no instabilities is to appeal to photon-axion mixing (conversion of photons to axions) in a universe dominated by a cosmological constant (or quintessence) [60].

In our models, the superacceleration arises due to interactions of dark energy and matter. Our models therefore provide super-acceleration with none of the atten-

dant problems that plague most of the above models. Furthermore, the interactions are generic; we do not need to fine-tune couplings in order to avoid theoretical pitfalls or observational constraints. We therefore believe that considering interactions of dark energy is the best way to generate models of superacceleration.

### AZK-INSTABILITY

In this section we will consider a general class of models in which the dark energy density is coupled to the non-relativistic matter density. For an example of how this could occur, suppose that non-relativistic matter particles are coupled to a scalar field. Thus the local density of the matter particles can influence the vacuum expectation value (vev) of the scalar field. The change in the potential of the scalar then affects the dark energy, thus coupling matter and dark energy.

In this class of models, the matter fields will be taken to have a matter density  $n_M$ . They are coupled to a scalar field  $\chi$  (dark energy) through Yukawa like couplings. We take the potential to be

$$E = \int d^3x V(\chi, n_M), \quad (1)$$

$$= \int d^3x [V_0(\chi) + mn_M + \lambda g(\chi)n_M]. \quad (2)$$

We will assume that  $m_\chi^2 = V_0''(\chi_0) + \lambda g''(\chi_0)n_M$ , the mass-squared of the scalar field about its vev  $\chi_0$ , is very large so that the  $\chi$  field always sits at the minimum of its effective potential. This is the central assumption of our paper. The mass will certainly have to be larger than the expansion rate of the universe to be consistent with this assumption. We will also assume that the mass is large enough to satisfy the constraints imposed by experiments that probe the strength of a fifth force.

In the absence of the last term, this is the potential energy of two decoupled fluids. The first term corresponds to a cosmological constant term (since we have assumed that the field  $\chi$  is always at the minimum). The second term is the energy density of a dark matter fluid with density  $n_M$  and particle mass  $m$ .

The last term couples these two fluids, and leads to interesting effects. In particular  $\chi_0$ , the value of the scalar field at its minimum is now found by solving the equation

$$V_0'(\chi_0) + \lambda g'(\chi_0)n_M = 0, \quad (3)$$

where  $V_0'$  and  $g'$  are derivatives of  $V_0$  and  $g$  with respect to  $\chi$ . Thus  $\chi_0$  is now a function of  $n_M$ .

We can make the dependence of  $\chi_0$  on  $n_M$  explicit in the following way. Consider small deviations in  $n_M$ . The vev of the scalar field shifts to account for this change in  $n_M$ . Taking a further derivative, we find

$$(V_0''(\chi_0) + \lambda g''(\chi_0)n_M) \frac{\partial \chi_0}{\partial n_M} + \lambda g'(\chi_0) = 0. \quad (4)$$

This explicitly shows how  $\chi_0$  varies as  $n_M$  varies.

In writing Eq. 2, we neglected the kinetic term in comparison to the potential. This is necessary if the scalar field is to behave as dark energy and, as we now show, consistent with our assumption of a large mass for the scalar field. Note that  $\dot{\chi} = n_M \partial \chi_0 / \partial n_M$ . Working out this expression, we find that  $\dot{\chi}^2/V$  for  $\chi = \chi_0$  is given by  $(V_0'^2/Vm_\chi^2)(n_M/m_\chi n_M)^2$ . Lets look at changes to the scalar field potential around  $\chi = \chi_0$ . Unless there are strong fine-tunings and cancellations, we will have  $V_0'\delta\chi < V$  and  $m_\chi^2(\delta\chi)^2/2 < V$ , which together imply that  $2V_0'^2/Vm_\chi^2 < 1$ . Hence the natural expectation is that  $\dot{\chi}^2/V \sim H^2/m_\chi^2$ . For large enough  $m_\chi$ , the kinetic term is negligible.

We now show that there is an instability in this system. We start with a configuration where the dark matter is evenly distributed, and the  $\chi$  field is at its minimum  $\chi_0$  everywhere. Now consider small fluctuations in the matter density  $\delta n_M$  which preserve  $\int_\tau d^3x \delta n_M = 0$ , i.e., the total number in volume  $\tau$ . The integral is over some region  $\tau$ , much smaller than the Hubble volume, over which the fluctuations are coherent. Such a fluctuation leads to a change in the total energy. The energy change proportional to  $\delta n_M$  vanishes because of Eq. 3 and the condition that  $\int_\tau d^3x \delta n_M = 0$ . The energy change to next order is

$$\delta E = \frac{1}{2} \int d^3x (\delta n_M)^2 \left( \frac{\partial \chi_0}{\partial n_M} \right) \left( m_\chi^2 \frac{\partial \chi_0}{\partial n_M} + 2\lambda g'(\chi_0) \right) \quad (5)$$

$$= -\frac{1}{2} \int d^3x (\delta n_M)^2 \lambda^2 \frac{[g'(\chi_0)]^2}{m_\chi^2} \quad (6)$$

Therefore the leading correction to the energy is always negative, implying that the configuration is unstable to the growth of these fluctuations. We dub this the AZK-instability. This instability was first noted in the context of the MaVaN scenario [12].

We have neglected gravity and the expansion of the universe in the above analysis. We neglected gravity because the relevant length scales are much smaller than the Jeans length; the instability occurs on all scales and hence the effect is most severe on microscopic scales. The analysis above was thus for a region  $\tau$  much smaller than that where gravity would be important. We neglected the expansion of the universe because the relevant time scales are much smaller than the age of the universe. In addition our setup started with a smooth distribution of matter. For this one must go to scales smaller than the free-streaming scale of dark matter particles. For example, the comoving free-streaming scale of a typical neutralino dark matter particle is of the order of parsec. We do not study this system on larger cosmologically relevant scales. It is, however, unlikely that the system will still be able to drive the accelerated expansion of the universe since the generic AZK instability is intimately related to the adiabatic sound speed of the fluid [12].

The result above assumes that the scalar field is much

heavier than the expansion rate of the universe. This constraint is easy to satisfy and the large mass makes the model more robust to radiative corrections (for example, see [16]). Secondly, the calculation is only valid for modes which have a wavelength much larger than  $1/m_\chi$ ; for shorter wavelengths, we cannot assume that the scalar field relaxes to the minimum quickly enough.

### AVOIDING THE AZK-INSTABILITY

To avoid this instability, we look at more general couplings.

Consider now a model where the total energy is

$$E = \int d^3x [V_0(\chi) + mn_M + \lambda g(\chi)n_M^n], \quad (7)$$

and we choose  $\lambda > 0$  without loss of generality.

Again we assume that the scalar field tracks the minimum of the potential and hence we have,

$$V'_0(\chi_0) + \lambda g'(\chi_0)n_M^n = 0, \quad (8)$$

$$(V''_0(\chi_0) + \lambda g''(\chi_0)n_M^n) \left( \frac{\partial \chi_0}{\partial n_M} \right) + \lambda g'(\chi_0)n_M^{n-1} = 0. \quad (9)$$

Following our earlier calculation, we find

$$\delta E = \frac{1}{2} \int d^3x \left( \frac{\delta n_M}{n_M} \right)^2 \left( -\frac{[n\lambda g'(\chi_0)n_M^n]^2}{m_\chi^2} + \lambda n(n-1)g(\chi_0)n_M^n \right), \quad (10)$$

Therefore, the instability is avoided if

$$-n^2\lambda^2 n_M^n \frac{[g'(\chi_0)]^2}{m_\chi^2} + n(n-1)\lambda g(\chi_0) > 0. \quad (11)$$

We note that the first term is always negative and gets large with  $n_M$  unless  $g'(\chi_0)$  decreases fast enough. Looking at the second term we note that any value of  $0 < n \leq 1$  is unstable independent of the form of  $g(\chi)$  except for the requirement that  $g(\chi_0) > 0$  which is required anyway for the potential to be bounded from below.

A robust way to avoid the instability is to choose  $n < 0$ , which makes the second term positive. This is, of course, not sufficient to guarantee the inequality in Eq. 11. We need the magnitude of the second term to be larger than that of the first. This is easy to arrange. We again look at changes to the potential as we vary  $\chi$  about  $\chi_0$ . If the potential is not fine-tuned to give rise to cancellations between terms in the Taylor expansion, then  $n\lambda g'n_M^n \delta\chi < V$  and also  $m_\chi^2(\delta\chi)^2/2 < V$ . Putting these two expressions together yields  $2n^2\lambda^2(g')^2 n_M^{2n}/m_\chi^2 < V \sim \lambda g n_M^n$ . Hence we see that it is natural, if  $n < 0$ , for the inequality in Eq. 11 to be satisfied.

It is also possible to avoid the instability by choosing  $n > 1$ . However, this region of model space will be heavily constrained by observations. In situations where the matter density gets large, i.e., in collapsed structures, the last term in the potential dominates. It would make the dark energy density in galaxies large, change structure formation and clustering properties of dark matter halos. Therefore, these kinds of models would be tightly constrained. In order for these models to be viable,  $\lambda$  would have to be small and the model would essentially be the same as that with two decoupled fluids.

Thus the requirement of AZK-stability and observational constraints naturally lead us to consider models where  $n < 0$ . We now look at observational consequences of such a coupling.

### AZK-STABILITY AND SUPER-ACCELERATION

The coupling term above with  $n < 0$  introduces a very interesting effect: this model has super-acceleration. That is, observations will seem to show a phase with dark energy equation of state less than -1.

To see this, we first note that the observational quantity that is important is the pressure. We will fit to the observations a model with matter that scales with the expansion as  $1/a^3$ , and dark energy with some equation of state  $w_{DE}$ . Note that adding or removing a component of energy density that scales as  $1/a^3$  does not change the pressure of the fluid. Hence very generally  $P_{tot} = P_{DE}$ .  $P_{tot}$  is defined by the equation  $\dot{V} = -3H(V + P_{tot})$  from which we find  $P_{tot} = -V_0(\chi_0) + \lambda g(\chi_0)n_M^n(n-1)$ . We set the equation of state  $w_{DE} \equiv P_{tot}/(V - mn_M)$  and find,

$$w_{DE} = -1 + \frac{n\lambda g(\chi_0)n_M^n}{V_0(\chi_0) + \lambda g(\chi)n_M^n}. \quad (12)$$

Now since  $n < 0$ , the second term is actually negative, and we have  $w_{DE} < -1$  i.e. super-acceleration.

We emphasize that this super-acceleration is *not* accompanied by any of the problems normally associated with theories with equation of state less than -1. There is no acausal behavior, and there are no ghosts. This is because the super-acceleration in our model results from an interaction which is ignored in the fitting of theory to observations. If we fit our observations using a canonical matter density term and dark energy, then the interaction has the effect of making the effective equation of state for dark energy more negative.

### SOUND SPEED

Here we present an alternative derivation of the instability in terms of the sound speed of the combined fluid. A negative sound speed squared would signal instability.

On length scales much larger than  $m_\chi^{-1}$ , the evolution of the system is adiabatic and hence the sound speed is

$$c_a^2 = \frac{\dot{P}_{\text{tot}}}{\dot{V}}. \quad (13)$$

The adiabatic sound speed in this theory can then be expressed as

$$c_a^2 = \frac{n_M \partial w_{\text{DE}} / \partial n_M + w_{\text{DE}}(1 + w_{\text{DE}})}{1 + w_{\text{DE}} + m n_M / (V - m n_M)}. \quad (14)$$

$$= \frac{n_M}{M} \left[ \frac{\partial^2 V(\chi_0, n_M)}{\partial n_M^2} - m_\chi^2 \left( \frac{\partial \chi_0}{\partial n_M} \right)^2 \right], \quad (15)$$

$$= \frac{n_M \partial w_{\text{tot}} / \partial n_M + w_{\text{tot}}(1 + w_{\text{tot}})}{1 + w_{\text{tot}}}, \quad (16)$$

where  $w_{\text{tot}} \equiv P_{\text{tot}}/V$  is the equation of state of the total fluid.

For a universe with an accelerating expansion  $w_{\text{tot}} < -1/3$ . For a wide class of models with  $w_{\text{tot}} < 0$  and either the  $n_M w'_{\text{tot}}$  term sub-dominant or negative, we have  $c_a^2 < 0$  and the system is unstable. This is just the AZK-instability.

Lets now look in more detail at Eq. 14. First, consider the case where  $w_{\text{DE}} > -1$ : the denominator is positive and if the  $w'_{\text{DE}}$  term is sub-dominant or negative, then AZK-instability sets in. It is clear that this instability may not be present in models with  $w_{\text{DE}} < -1$ . We also note that this instability will likely set in well before the current epoch because at early times  $n_M/(V - n_M) \gg 1$ . For this case where  $w_{\text{DE}}(1 + w_{\text{DE}}) > 0$ , the sign and magnitude of the  $n_M w'_{\text{DE}}$  term is important. In particular, the requirement that the  $n_M w'_{\text{DE}}$  term is sub-dominant may not be trivial to obtain [61].

While the above derivation shows us how the instability arises, it does not provide us with an intuitive understanding of what happens to the matter. In order to better understand that we look at the Boltzmann equation for the matter coupled to a scalar field. The scalar field gives the matter a mass term that can vary spatially and temporally. Following AZK [12], we write down the Boltzmann equation for matter neglecting gravity and hence only valid on small scales. These are the scales of interest since we have assumed  $m_\chi \gg H$ . We write down the first order perturbations to this equation and expand the perturbations in plane wave modes. Denoting the effective mass of the matter particle by  $M(\chi)$  we find,

$$\omega \delta f(\mathbf{p}, \mathbf{k}) - (\gamma M)^{-1} \mathbf{p} \cdot \mathbf{k} \delta f(\mathbf{p}, \mathbf{k}) \quad (17)$$

$$+ \gamma^{-1} \delta M(\mathbf{k}) \mathbf{k} \cdot \nabla_{\mathbf{p}} f(\mathbf{p}) = 0. \quad (18)$$

We then find the perturbation to the matter density  $\delta n_M(\mathbf{k})$  using the above equation. In the limit that matter is non-relativistic, the resulting equation has a simple form. We find that the variation in effective mass of the particle is given by  $\delta M(\mathbf{k}) = (M/n_M) c_s^2 \delta n_M(\mathbf{k})$  where we

have defined  $c_s = \omega/k$ , the sound speed of matter. The above equation is valid for perturbations  $\delta M$  on all scales at which our assumptions hold. As pointed out in [12], there is no scale in the equation for  $c_s^2$  because we are studying scales where it is correct to assume that the scalar field adjusts to changes in the matter density, and gravity is unimportant.

We now turn to the fluid description and write  $M = V(\chi_0, n_M)/\partial n_M$ . Using Eq. 4 for  $d\chi_0/dn_M$ , one may then obtain perturbations in  $M$  as  $\delta M = (M/n_M) c_a^2 \delta n_M$  where  $c_a^2$  is given by Eq. 15. In the framework of a scalar degree of freedom coupled to matter, both descriptions must be valid and hence we find that  $c_s^2 = c_a^2$ . The instability may therefore be analyzed in terms of  $c_a^2$ . All of our analyses in earlier sections go through if we work with  $c_a^2$  and we conclude that models with super-acceleration provide a generic way to avoid the AZK instability.

## CONCLUSIONS

In this paper, we have explored the possibility that dark energy may interact with matter. Such a hypothesis is natural if the explanation for dark energy requires extra scalar degrees of freedom. Unfortunately, as we have shown here, these models suffer from a generic instability when the mass of the scalar field is very large. We have verified that this instability is also present in scalar-tensor theories where the scalar plays the role of dark energy, and also in models with multiple scalar fields.

We then looked for models where this instability could be avoided, and found a large class of such models. Most interestingly, we found that in these models, the *apparent* equation of state of the dark energy density is generically smaller than -1. This super-acceleration is a result of the fact that we fit observations with models that have non-interacting matter and dark energy fluids.

There is a theoretical prejudice against models of  $w_{\text{DE}} < -1$  due to their apparent theoretical problems. The observational data certainly do not disfavor  $w_{\text{DE}} < -1$ . Indeed a large region of the parameter space allowed by SNIa observations corresponds to a constant  $w_{\text{DE}} < -1$ . Here we have shown that stable models with  $w_{\text{DE}} < -1$  may be constructed without encountering ghosts or acausal behavior. These models are no more fine-tuned than quintessence models. Thus theoretical bias against  $w_{\text{DE}} < -1$  should be treated with circumspection, and not be given any weight when interpreting observational data.

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